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We present a model with minimal assumptions for non-thermal leptogenesis with almost degenerate superheavy right-handed neutrinos in a supersymmetric set up. In this scenario a gauge singlet inflaton is directly coupled to the right-handed (s)neutrinos with a mass heavier than the inflaton mass. This helps avoiding potential problems which can naturally arise otherwise. The inflaton decay to the Standard Model leptons and Higgs, via off-shell right-handed (s)neutrinos, reheats the Universe. The same channel is also responsible for generating the lepton asymmetry, thus requiring no stage of preheating in order to excite superheavy right-handed (s)neutrinos. The suppressed decay rate of the inflaton naturally leads to a sufficiently low reheat temperature, which in addition, prevents any wash out of the yielded asymmetry. Finally, a successful leptogenesis can be accommodated for a variety of inflationary models with a rather wide ranging inflationary scale.

I. INTRODUCTION

The consistency of the abundance of the light elements synthesized during the big bang nucleosynthesis (BBN) requires that the baryon asymmetry of the Universe (BAU), parameterized as $\eta_B = (n_B - n_{\bar{B}})/s$ with s being the entropy density and n_B the number density of the baryons, be in the range $(0.3 - 0.9) \times 10^{-10}$ [1]. The asymmetry can be produced from a baryon symmetric Universe, provided three conditions are simultaneously met; B and/or L -violation, C - and CP -violation, and departure from thermal equilibrium [2]. However, any produced asymmetry will be washed away by the SM $B + L$ -violating sphaleron transitions which are active from temperatures 10^{12} GeV down to 100 GeV [3], if $B - L = 0$. Therefore, an asymmetry in $B - L$ is generally sought which is subsequently reprocessed in a thermal bath via sphalerons in order to yield a net baryon asymmetry given by $B = a(B - L)$. Here, a is a model-dependent parameter; in case of the standard model (SM), $a = 28/79$, while in the minimal supersymmetric standard model (MSSM), $a = 32/92$ [4].

An attractive mechanism for producing the $B - L$ asymmetry is from the decay of the heavy right-handed (RH) Majorana neutrinos [6]. Since the RH neutrinos are the SM singlets, a Majorana mass M_N , which violates lepton number is compatible with all symmetries and hence can be arbitrarily large beyond the electroweak scale. This provides an elegant way for obtaining small masses m_ν for the light neutrinos via the see-saw mechanism such that $m_\nu \approx (m_D^2/M_N)$ [5], where m_D is the Dirac mass obtained from the Higgs vacuum expectation value (VEV). Moreover, a lepton asymmetry can be generated from the interference between the tree-level and the one-loop diagrams in out-of-equilibrium decay of the RH neutrinos, provided CP -violating phases exist in neutrino Yukawa couplings. The lepton asymmetry thus obtained, will be partially converted to the baryon asymmetry via sphaleron effects. This is the standard

lore of producing lepton asymmetry commonly known as leptogenesis [6,7].

The present analyses of solar neutrino experiments favor the large mixing angle MSW solution with $\Delta m_{\nu, \text{solar}}^2 = 6.1 \times 10^{-5} \text{ eV}^2$ and $\tan^2 \theta_{12} = 0.41$ [8], while $\Delta m_{\nu, \text{atm}}^2 = 3.2 \times 10^{-3} \text{ eV}^2$ and $\sin^2(2\theta_{23}) = (0.83 - 1)$ provides the best fit to the atmospheric neutrino data [9].

In addition, cosmology [10], and neutrinoless double-beta decay experiments [11] provide an upper limit for the light neutrino masses. The masses and mixing angles which are required to explain solar and atmospheric neutrino data can be obtained in both scenarios with hierarchical, or quasi-degenerate neutrinos. Note that the hierarchical spectrum of heavy neutrinos strongly suggests a spectrum of light neutrinos which is hierarchical too, unless there is a big conspiracy. On the other hand, a mild hierarchy of RH neutrino masses could be compatible with degenerate light neutrinos with a certain amount of fine-tuning. In the former case, one may consider thermal leptogenesis scenario where heavy neutrinos come into equilibrium with the primordial thermal bath through Yukawa interactions. The decay of the lightest RH neutrino easily satisfies the out of equilibrium condition by virtue of having a sufficiently small Yukawa coupling [7]. In a model-independent analysis in Ref. [12], the authors have parameterized thermal leptogenesis by four parameters; the CP asymmetry, the heavy RH neutrino mass, the effective light neutrino mass, and the quadratic mean of the light neutrino masses. The final result was that an acceptable lepton asymmetry can be generated with $T_R \sim M_1 = \mathcal{O}(10^{10})$ GeV, and $\sum_i m_{\nu, i} < \sqrt{3} \text{ eV}$.

However, the temperature required for thermal leptogenesis is marginally compatible with the maximum allowed one in supersymmetric theories, which is usually constrained by thermal gravitino production [13,14]. Gravitinos with a mass $\mathcal{O}(\text{TeV})$ decay after nucleosynthesis and their decay products can change abundance of the light elements synthesized during BBN. For $100 \text{ GeV} \leq$

$m_{3/2} \leq 1$ TeV, a successful nucleosynthesis requires $n_{3/2}/s \leq (10^{-14} - 10^{-12})$, which translates into $T_R \leq (10^7 - 10^{10})$ GeV [13,14] *. The possible ways for obtaining a naturally low reheat temperature include gravitationally suppressed decay of the inflaton in models of high scale inflation [22], low scale inflationary models [23–25], a brief period of late thermal inflation [26], or, a completely new paradigm “reheating through the surface evaporation” which works even for high scale inflationary models [27].

When the light neutrinos are almost degenerate $m_{\nu,1} \approx m_{\nu,2} \approx m_{\nu,3}$, which requires quasi-degenerate heavy neutrinos, the out of equilibrium condition in thermal leptogenesis scenario cannot be satisfied in the minimal see-saw model [7]. More complicated models are required in this case [28]. On the other hand, if the mass splitting of the RH neutrinos becomes less than their decay widths, the perturbative calculations obviously break down. Then, the effect of finite decay widths of the RH neutrinos must be taken into account [29]. The careful treatment of Ref. [29] shows that a resonant enhancement of lepton asymmetry occurs in this case, while as expected, it vanishes in the limit of exactly degenerate neutrinos. This effect can be utilized to bring down the scale of heavy neutrino masses, and hence the leptogenesis scale [30].

However, for almost degenerate heavy neutrinos, i.e. where the mass splitting is larger than the decay width, one has to seek non-thermal leptogenesis (which works for the hierarchical neutrino masses as well) in the minimal models. In this scenario RH neutrinos are produced non-thermally from the inflaton decay. This can occur during reheating if the inflaton decays to the RH neutrinos, which are lighter than the inflaton, with a considerable branching ratio [31]. Heavy neutrinos can also be produced via preheating [32] (a stage of reheating where a resonant production of massive and/or massless bosons and fermions takes place [33]) or tachyonic preheating [34], even if the mass of boson/fermion exceeds that of the inflaton. However, this is rather model-dependent and its main features can significantly vary from model to model. This is the prime reason why we do

not pursue leptogenesis via preheating mechanism here. In supersymmetric models one has the RH sneutrinos in addition. The sneutrinos are produced along-with neutrinos during reheating, and with much higher abundances in preheating, thus serving as an additional source for leptogenesis [35]. Moreover, the RH sneutrinos can acquire a large VEV during inflation, if their mass is less than the Hubble expansion rate during inflation H_I . Such a condensate starts oscillating once $H(t) \simeq M_N$, therefore, automatically satisfying the out of equilibrium condition. The decay of the sneutrino condensate can then yield the desired lepton asymmetry in the same fashion as neutrino decay [36], or via Affleck-Dine mechanism [37]. This last scenario has an additional advantage that it solves the fine-tuning problem in F -term hybrid inflationary model in a very natural way [38].

The success of all these scenarios, but preheating and the Affleck-Dine oriented model, requires that the inflaton is heavier than the RH (s)neutrinos (in the hierarchical case inflaton only needs to be heavier than the lightest RH (s)neutrino.). Moreover, all the above scenarios are based upon the decay processes. An attractive proposal was recently made, where the lepton asymmetry in the visible sector is generated from the RH neutrino-mediated scattering of the SM Higgs and leptons into a depleted hidden sector [39], rather than the decay of the on-shell heavy neutrinos.

In this note we propose a simple supersymmetric model for non-thermal leptogenesis *without* any need of preheating mechanism. In this model the inflaton is directly coupled to nearly degenerate RH (s)neutrinos which are heavier than the inflaton. Then the inflaton decay to the SM fields, occurring mainly via off-shell RH (s)neutrinos, reheats the Universe and naturally leads to a sufficiently low reheat temperature. This same channel is also responsible for producing the lepton asymmetry.

In the next section we introduce our model and highlight several of its advantages. Then we turn to reheating and generation of the lepton asymmetry in this model and present our main results. Finally, we conclude the paper with a brief summary.

II. THE MODEL

We start by introducing our model in a supersymmetric set up. The relevant part of the superpotential is given by

$$W \supset \frac{1}{2}m_\phi\Phi\Phi + \frac{1}{2}g\Phi\mathbf{N}\mathbf{N} + h\mathbf{N}\mathbf{H}_u\mathbf{L} + \frac{1}{2}M_N\mathbf{N}\mathbf{N}. \quad (1)$$

Here Φ , \mathbf{N} , \mathbf{L} , and \mathbf{H} stand for the inflaton, the RH neutrino, the lepton doublet, and the Higgs (which gives mass to the top quark) superfields, respectively. Also, m_ϕ and M_N denote inflaton and RH (s)neutrino masses,

*Recently, non-thermal production of helicity $\pm 3/2$ [15], and helicity $\pm 1/2$ gravitinos [16,17] from inflaton oscillations have been considered. For a single chiral multiplet the helicity $\pm 1/2$ gravitino is the superpartner of the inflaton known as inflatino. The decay channels of inflatino have been discussed in Ref. [18]. Also, it has been suggested [18], and explicitly shown [19], that in realistic models with two chiral multiplets the helicity $\pm 1/2$ gravitino production is not a problem, so long as the inflationary scale is sufficiently higher than the scale of supersymmetry breaking in the hidden sector and the two sectors are weakly coupled. Gravitinos can also be produced directly from the inflaton decay [20], and from the decay of the heavy stable neutral particles [21].

respectively[†]. We assume that the inflaton is coupled to the RH (s)neutrinos via Yukawa coupling g , and h denotes the neutrino Dirac Yukawa matrix. For simplicity, we have omitted all indices in h matrix and superfields, and work in the basis where M_N is diagonal. Further simplifications can be made for almost degenerate RH (s)neutrinos where M_N is essentially the same for all of them. It is also conceivable in this case that the inflaton is coupled with the same strength to three RH (s)neutrinos. This is particularly true when the inflaton has a nonzero VEV at the minimum which provides masses to the RH (s)neutrinos. We focus on superheavy RH (s)neutrinos, i.e. assuming that $M_N \gg m_\phi$.

Now let us discuss the merits why we seek RH (s)neutrinos heavier than the inflaton. If $M_N < m_\phi$, then one can easily produce on-shell (s)neutrinos from the inflaton decay, either perturbatively or via preheating. First consider (s)neutrino production in perturbative inflaton decay. A perturbative decay requires a small coupling to the (s)neutrinos. This is naturally achieved when the inflaton lies in a hidden sector which is only gravitationally coupled to the SM sector [22]. In this case the total decay rate of the inflaton is given by $\Gamma_d \sim m_\phi^3/M_P^2$, whilst the partial decay rate to (s)neutrinos is given by $\Gamma_{\phi \rightarrow N} \sim m_\phi M_N^2/M_P^2$ [21], where $M_P = 2.4 \times 10^{18}$ GeV is the reduced Planck mass. This results in a branching ratio $\simeq (M_N/m_\phi)^2$, which is too small if $M_N \ll m_\phi$. Note that a successful leptogenesis requires an acceptable branching ratio given the entropy generation from reheating, thus implying that M_N must not be much smaller than m_ϕ .

Besides, a small coupling g , which is required to ensure a perturbative treatment, leads to another potential problem. The sneutrino field \tilde{N} can acquire a large VEV during inflation. Since \tilde{N} is directly coupled to the inflaton, it might even ruin the flatness of the inflaton potential. On a lighter note, $\langle \tilde{N} \rangle$ remains non-vanishing after the end of inflation in any case, and may contribute to isocurvature density perturbations [41]. This requires a delicate treatment of a coupled system which depends on the choice of a model. This is an issue which has been side-lined in most supersymmetric models of non-thermal leptogenesis, except Ref. [37].

If g is sufficiently large, (s)neutrinos may be produced in a non-perturbative manner during the stage of (tachyonic) preheating [32,34]. For the superpotential in Eq. (1), the necessary condition for preheating reads $g\phi_0 > m_\phi$, where ϕ_0 is the initial amplitude of the inflaton oscillations. This guarantees that \tilde{N} is heavier than the inflaton during inflation, and hence $\langle \tilde{N} \rangle = 0$

after the end of inflation, resulting in a simpler initial condition in post-inflationary era. On the other hand, both RH neutrinos and sneutrinos can be produced via preheating (sneutrinos much more abundantly by virtue of obeying the Bose statistics [33]). However, as mentioned earlier, this is rather model-dependent. For example, if the inflaton has a VEV at the minimum, denoted as v , then it is hard to envisage an efficient production of (s)neutrinos through parametric resonance. The reason is that $M_N = gv$ and $\phi_0 \simeq v$ in this case, which implies $g\phi_0 \simeq M_N$. It is therefore evident that there will be no preheating of (s)neutrinos for $M_N < m_\phi$. On the other hand, for $M_N > m_\phi$ preheating is possible only if $g\phi_0 m_\phi \gg M_N^2$ [33,42]. In particular, resonant creation rapidly ceases to be efficient for $M_N > 10m_\phi$ [42][‡]. In the tachyonic preheating scenario, too, the produced (s)neutrinos usually have an abundance much less than the inflaton abundance when $M_N \gg m_\phi$ [34]. In conclusion, it is very difficult (if not impossible) to obtain the desired lepton asymmetry in a wide range of inflationary models, by solely relying on non-perturbative dynamics.

Now we count upon the advantages of our model. First of all, for $M_N \gg m_\phi$ the post-inflationary dynamics is simpler since $\langle \tilde{N} \rangle = 0$ at the end of inflation. The Universe is reheated through the inflaton decay to the Higgs and SM leptons via off-shell RH (s)neutrino. The decay rate, as we will see shortly, is suppressed as $(m_\phi/M_N)^4$. This naturally leads to an acceptably low reheat temperature when $M_N \gg m_\phi$. Furthermore, the inflaton decay alone is responsible for the generation of the lepton asymmetry. This makes the model minimal since leptogenesis is now directly connected with reheating. Also, the washing out of the lepton asymmetry from thermal scatterings of the SM leptons and Higgs is completely negligible since $T_R \ll M_N$.

Our main focus will be on almost degenerate light neutrinos, which can be derived naturally from almost degenerate RH neutrinos. An example of such a model is presented in Ref. [44], where neutrino masses and mixings compatible with the solar and atmospheric neutrino solutions are derived in the framework of democratic mass matrix. There, the neutrino Yukawa matrix h is almost diagonal in the same basis as M_N , i.e. $h_{ij}, \Delta h_{ii} \ll h_{ii}$. This makes sense since when both M_N and h are proportional to the identity matrix, the light neutrinos come out to be exactly degenerate. Then, by perturbing M_N and h around this pattern, we can obtain a nearly degenerate texture. In the calculations below, M_N and h

[†]Actually, m_ϕ denotes the frequency of the inflaton oscillations around the global minimum of the potential. In models of chaotic inflation $m_\phi \simeq H_I$. While in new and hybrid inflationary models it is usually (much) larger than H_I [40].

[‡]It has been shown in Ref. [42], that for a quadratic potential; $V_\phi \sim m_\phi^2 \phi^2$, efficient resonant production of particles with a mass $M_N = 10m_\phi$ requires $g\phi_0 > 10^4 m_\phi$. On the other hand, for a quartic potential, $V_\phi \sim \lambda \phi^4$, preheating of these particles practically disappears. Preheating in supersymmetric hybrid inflation model is also not efficient [43]

denote nearly equal diagonal elements of the Majorana mass and neutrino Yukawa matrices, respectively.

III. REHEATING THE UNIVERSE

The main decay mode of the inflaton is to a four-body final state consisting of two Higgs/Higgsino-lepton/slepton particles. Since, we have assumed $m_\phi \ll M_N$, it is essential to find those diagrams which are least suppressed by powers of M_N . These diagrams, shown in Fig. (1), arise from the leading order terms in the effective superpotential which, after integrating out \mathbf{N} , is given by

$$W_{eff} \supset \frac{1}{2} m_\phi \Phi^2 + \frac{1}{2M_N^2} g h^2 \Phi (\mathbf{H}_u \mathbf{L}) (\mathbf{H}_u \mathbf{L}). \quad (2)$$

We should therefore choose that part of the N propagator with a mass insertion, namely the part suppressed as $1/M_N$ (the other part of the propagator is proportional to m_ϕ/M_N^2). In diagrams below, two opposite arrows on the N propagator represent this dominant part. Note that the \tilde{N} propagator is proportional to $(1/M_N^2)$ to the leading order.

Generally, the trajectory of inflaton motion is a line in the complex ϕ plane. We can therefore assume, without loss of generality, that only the real component of the inflaton has a VEV, thus treating the decaying inflatons as real fields. In addition, the SM particles are much lighter than the inflaton in the case under consideration (as will be confirmed by our results). Then the phase space factor for the four-body decay is readily found to be $[16 \cdot 96 \cdot (2\pi)^5]^{-1}$. There is also an additional factor of three, since the decay rate receives the same contribution from each N (\tilde{N}), recall that each N (\tilde{N}) is dominantly coupled to only one of the lepton doublets. After collecting all factors of g and h and summing up over all possible final states, with different contributions to the same final state properly taken into account, we obtain the total decay rate of the inflaton:

$$\Gamma_d \simeq \frac{17}{2^{14}\pi^5} g^2 h^4 \frac{m_\phi^5}{M_N^4}. \quad (3)$$

The inflaton completely decays when $H \simeq \Gamma_d$, where $H \simeq (g_*^{1/2} T^2/M_P)$ in a radiation-dominated Universe [40], with g_* being the effective number of relativistic degrees of freedom which is $\simeq 214$ in the MSSM. Assuming that thermal equilibrium is achieved when $H \simeq \Gamma_d$ (this is justifiable, for the detailed discussion see Refs. [45,46]), we obtain

$$\frac{T_R}{m_\phi} \simeq 10^{-7/2} \frac{g h^2 m_\phi^{3/2} M_P^{1/2}}{M_N^2}. \quad (4)$$

Some comments are in order regarding our estimates of Γ_d and T_R . One might think that inflaton decay into four scalars, the same as in diagram (b) except that \tilde{H}_u

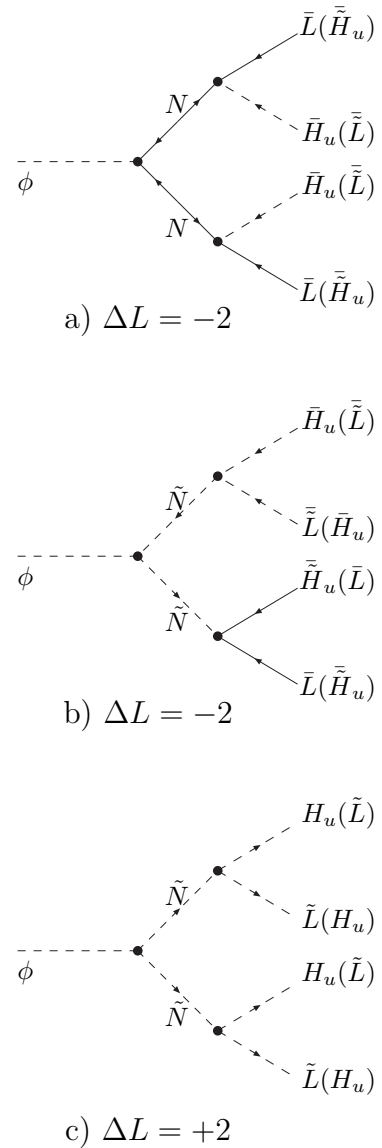


Fig. 1: Diagrams which, together with their CP transformed, represent inflaton decay into two Higgs/Higgsino-lepton/slepton pairs at the leading order.

and \tilde{L} are replaced with H_u and \tilde{L} , would occur at a rate only suppressed by two powers of M_N . However, this is not the case since this leading order contribution is canceled out by that from another diagram and the overall rate is actually proportional to (m_ϕ^7/M_N^6) . This is just a manifestation of the fact that these diagrams do not arise from the effective superpotential individually. Also, there exists a two-body decay channel for the inflaton, into two $\tilde{H}_u H_u(\tilde{H}_u \tilde{H}_u)$ or $\tilde{L} L(\tilde{L} \tilde{L})$, at the one-loop level. It can easily be derived by choosing $(1/M_N)$

and (m_ϕ/M_N^2) parts of N propagators in diagram (a) and connecting the $\bar{H}_u(\bar{L})$, or $\bar{L}(\bar{H}_u)$, lines. This channel has a much larger phase space factor $8\pi^{-1}$, while the dependence on g and h remains the same as in Fig. (1). However, the two-body decay rate is $\propto (m_\phi^7/M_N^6)$. Thus, by taking the one-loop factor $4\pi^{-2}$ into account and for $M_N \geq 10m_\phi$, it will eventually be smaller than that in Eq. (3).

Finally, the inflaton can also decay into the SM fields via gravitational couplings with a decay rate $\Gamma_{grav} \sim (v/M_P)^2(m_\phi^3/M_P^2)$, where v denotes inflaton VEV at the global minimum of the potential [21]. Such a decay rate can however be neglected compared to the four-body decay provided $v \ll M_P$.

IV. THE LEPTON ASYMMETRY

In this section we evaluate the lepton asymmetry generated from the inflaton decay through diagrams in Fig. (1). A net lepton asymmetry is generated from the interference between these diagrams, and the one-loop diagrams representing self-energy and vertex corrections to one of the N (\tilde{N}) propagators. Here, we remind the readers that for the standard case where the decay of on-shell neutrinos yields the lepton symmetry, one has $\eta_L = \Sigma_i \epsilon_i (n_{N_i}/s)$, where

$$\epsilon_i = -\frac{1}{8\pi} \frac{1}{[hh^\dagger]_{ii}} \text{Im} \sum_{i \neq j} ([hh^\dagger]_{ij})^2 f\left(\frac{M_j^2}{M_i^2}\right), \quad (5)$$

and [7]

$$f(x) = \sqrt{x} \left(\frac{2}{x-1} + \ln \left[\frac{1+x}{x} \right] \right). \quad (6)$$

The first and second terms on the right-hand side of Eq. (6) correspond to the one-loop self-energy, and vertex corrections, respectively. It is easy to see that for almost degenerate N , the self-energy contribution dominates. Moreover, it can be shown that [47]

$$|\epsilon_1| \simeq |\epsilon_2| \simeq |\epsilon_3| \simeq \frac{3}{8\pi} \frac{M_N}{\langle H_u^0 \rangle^2} \frac{\Delta m_{\nu, atm}^2}{2m_\nu}. \quad (7)$$

Here $\langle H_u^0 \rangle = 174 \sin\beta$ GeV is the VEV of H_u in our vacuum, where $\tan\beta$ is defined as the ratio of $\langle H_u^0 \rangle$ and $\langle H_d^0 \rangle$.

In our case, where the lepton asymmetry is generated from off-shell N (\tilde{N}), the tree-level and one-loop diagrams are proportional to $1/M_N^2$ and $1/M_N^4$, respectively. Therefore ϵ_i get an extra suppression factor of $(m_\phi/M_N)^2$. More precisely, the suppression will go as $(m_\phi/2M_N)^2$ since the momentum flowing in each N (\tilde{N}) propagator is $m_\phi/2$ on average. However, for almost degenerate neutrinos with $\Delta M_N \leq 0.1M_N$, the factor from self-energy correction, in Eq. (6), will compensate this

by more than enough. Also, the leading order contributions from the three RH (s)neutrinos will add constructively [44]. Since, the self-energy correction can come from either of N (\tilde{N}) propagators, and the lepton number is violated by two units in inflaton decay, the total asymmetry in the baryons (after taking into account of sphaleron effects) can be expressed as[§]

$$\begin{aligned} \eta_B &= \left(\frac{n_B}{n_\phi} \right) \left(\frac{n_\phi}{s} \right) \\ &\simeq \frac{3}{2\pi} \left(\frac{M_N}{\langle H_u^0 \rangle^2} \right) \left(\frac{\Delta m_{\nu, atm}^2}{2m_\nu} \right) \times \\ &\quad \left(\frac{m_\phi}{M_N} \right)^2 \left(\frac{T_R}{m_\phi} \right), \end{aligned} \quad (8)$$

where $s = (2\pi^2/45)g_*T_R^3$. Note that n_ϕ/s denotes the dilution from reheating. By using Eq. (4), and the relationship $m_\nu \simeq (h^2 \langle H_u^0 \rangle^2 / M_N)$, we eventually obtain

$$\eta_B \simeq 10^{-49/2} g \frac{m_\phi^{7/2} M_P^{1/2}}{M_N^2 \langle H_u^0 \rangle^4} (1\text{GeV})^2, \quad (9)$$

where we have taken $m_\nu \approx 0.1$ eV. We also use $\langle H_u^0 \rangle = 174$ GeV in below.

Let us now present some numerical examples for $M_N \geq 10m_\phi$. This choice is compatible with the assumption of superheavy RH (s)neutrinos. With $M_N = 10m_\phi$ and $g = 1$ (10^{-3}), the desired baryon asymmetry (approximately one part in 10^{10}) is obtained for $m_\phi = 10^{11}$ (10^{13}) GeV, and $T_R = 10^7$ GeV. Note that there is no gravitino problem for such a low reheat temperature.

The merits of our model are already evident from these numbers. First of all, $T_R \ll M_N$ guarantees that lepton number violating scattering of the SM particles is completely negligible. Especially, keeping in mind that in the MSSM, there are larger number of scattering processes which can considerably attenuate the obtained asymmetry if the reheat temperature T_R is close to M_N [7]. In our case, obtaining sufficiently low reheat temperature is more than welcome in this regard. Also, the robustness of the inflaton mass m_ϕ lies in a range which favors both high and intermediate inflationary models, thus making the scenario more flexible. Note that by increasing the ratio M_N/m_ϕ , successful leptogenesis requires a larger inflaton mass for the same g , or a larger g for the same m_ϕ .

Now we briefly touch upon the case of degenerate neutrinos, where the mass splitting is less than the decay rate of the RH (s)neutrinos. In this case it is possible to obtain a resonant enhancement of the lepton, and hence

[§]The asymmetries yielded in each of the $\Delta L = -2$ and $\Delta L = +2$ decays are equal, since the Yukawa couplings appearing in the latter are the complex conjugate of those in the former.

baryon, asymmetry [29]. This will further add to the robustness of the model. For example, an enhancement of n_B/n_ϕ by a factor of 1000 implies that an acceptable asymmetry can be produced with $m_\phi = 10^9$ (10^{11}) GeV, for $g = 1$ (10^{-3}), and with $M_N = 10m_\phi$. Therefore, in principle it is possible to further lower the scale of inflation, while still obtaining the desired asymmetry.

And one final note before closing this section. A small number of on-shell (s)neutrinos might also have been produced non-perturbatively, from an inefficient preheating, and hence contribute to the resultant asymmetry through their decay. The asymmetry yielded in the decay of on-shell particles, denoted as η_B^{on} , will be

$$\eta_B^{on} \simeq \frac{M_N^2}{m_\phi^2} \left(\frac{n_{\tilde{N}} + n_N}{n_\phi} \right) \eta_B. \quad (10)$$

Note that η_B^{on} does not contain the suppression factor $(m_\phi/2M_N)^2$. On the other hand, a factor of 4 will be lost, relative to the off-shell case, since the one-particle decays of on-shell \tilde{N} and N violate the lepton number by one unit. Thus, the contribution from possible on-shell (s)neutrinos can be neglected, provided $(n_{\tilde{N}} + n_N) < (m_\phi/M_N)^2 n_\phi$. For the range of parameters in above, this is generally the case [32,34,42].

V. CONCLUSION

In this note we have provided a simple example for non-thermal leptogenesis with nearly degenerate superheavy RH neutrinos in a supersymmetric set up. We assumed that the inflaton is lighter than the RH (s)neutrinos, thus naturally avoiding some potential problems which can naturally arise. The inflaton decay via off-shell (s)neutrinos reheats the Universe and the model is minimal in the sense that the same channel is also responsible for generating the lepton asymmetry. As usual, the asymmetry arises from the interference between the tree-level and the one-loop diagrams representing self-energy correction of (s)neutrinos, although off-shell in our case, provided neutrino Yukawa couplings contain CP -violating phases.

Finally, we briefly emphasize on remarkable advantages of this model. First of all, leptogenesis can be accommodated rather simply without relying on non-perturbative production of RH (s)neutrinos. It is particularly attractive that the desired baryon asymmetry can be directly generated in the final stage of reheating, which is perturbative, regardless of any model-dependent effects which might have resulted in a first stage of non-perturbative reheating. Secondly, the suppressed decay of the inflaton naturally leads to an acceptably low reheat temperature, which is compatible with the gravitino bound and also prevents any wash out of the yielded asymmetry. Also, the desired lepton asymmetry can be generated for a range of inflaton mass accessible in large

and intermediate scale models of inflation. A resonant enhancement of the asymmetry for degenerate neutrinos further strengthens this scheme.

Qualitatively, we expect that the non-thermal scenario also work in the case of hierarchical RH (s)neutrinos. However, a more careful study should be performed in order to compare the quantitative results with those obtained here.

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